Investigation of the Field-Tuned Quantum Critical Point in $CeCoIn_5$

V.R. Shaginyan *

[a] *Petersburg Nuclear Physics Institute, Gatchina, 188300, Russia*

Abstract

The main properties and the type of the field-tuned quantum critical point in the heavy-fermion metal $CeCoIn_5$ arisen upon applying magnetic fields B are considered within the scenario based on the fermion condensation quantum phase transition. We analyze the behavior of the effective mass, resistivity, specific heat, charge and heat transport as functions of applied magnetic fields B and show that in the Landau Fermi liquid regime these quantities demonstrate the critical behavior which is scaled by the critical behavior of the effective mass. We show that in the high-field non-Fermi liquid regime, the effective mass exhibits very specific behavior, $M^* \sim T^{-2/3}$, and the resistivity demonstrates the $T^{2/3}$ dependence. Finally, at elevated temperatures, it changes to $M^* \sim T^{-1/2}$, while the resistivity becomes linear in T. In zero magnetic field, the effective mass is controlled by temperature T, and the resistivity is also linear in T. The obtained results are in good agreement with recent experimental facts.

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*E–mail: vrshag@thd.pnpi.spb.ru

Magnetic-field tuning of quantum critical points (QCPs) in heavy-fermion (HF) metals becomes a subject of intense current interest because, as it is widely accepted, an understating of quantum criticality can clear up a mystery of fundamental physics of strongly correlated systems [1]. A fundamental question is whether the QCPs observed in HF metals are different and related to different quantum phase transition or their nature can be captured by the physics of a single quantum phase transition. To answer this question, we have at least to explore a particular quantum critical point in order to identify its nature. It can hardly be done on pure theoretical grounds since there can exist a great diversity of quantum phase transitions and corresponding QCPs in nature [2,3]. Therefore, mutually complementary experimental facts related to the critical behavior and collected in measurements on the same HF metal are of crucial importance for understanding the physics of HF metals. Obviously, such HF metal is to exhibit the critical behavior and has no additional phase transitions. For example, the HF metal $CeRu_2Si_2$ can be regarded as fit for such study because the measurements have shown neither evidence of the magnetic ordering, superconductivity nor conventional Landau Fermi-liquid (LFL) behavior down to ultralow temperatures [4]. Unfortunately, by now only precise ac susceptibility and static magnetization measurements at small magnetic fields and ultralow temperatures are known [4]. While additional measurements of such properties as the heat and charge transport and the specific heat could produce valuable information about the existence of Landau quasiparticles and their degradation and clarify the role of the critical fluctuations near the corresponding QCP. Such measurements on the HF metal $CeCoIn_5$ were recently reported [5–8]. It was shown that the resistivity $\rho(T)$ of $CeCoIn_5$ as a function of temperature T is linear in T in the absence of magnetic field [5]. Due to the existence of magnetic field-tuned QCP with a critical field $B_{c0} \simeq 5.1$ T, the LFL behavior is restored at magnetic fields $B \geq B_{c0}$ [6–8]. At the LFL regime, the measurements of the specific heat and the coefficient A in the resistivity, $\rho(T) = \rho_0 + A(B)T^2$, describing the electron-electron scattering, have demonstrated that the Kadowaki-Woods ratio, $K = A(B)/\gamma^2(B)$ [9], is conserved [7]. Here $\gamma(B) = C/T$, and C is the specific heat. It was also shown that the coefficient A diverges as $A(B) \propto (B - B_{c0})^{-\alpha}$, with $\alpha \simeq 4/3$ [6,8]. Moreover, a recent study of $CeCoIn_5$ in magnetic fields $B > B_{c0}$ have revealed that the coefficients $A(B)$ and $C(B)$, with $C(B)$ describing a T^2 contribution to thermal resistivity κ_r, possess the same critical field dependence $A(B) \propto C(B) \propto (B - B_{c0})^{-\alpha}$, so that the ratio $A(B)/C(B) = c$ [8]. Here c is a field-independent constant characterizing electron-electron scattering in metals and having a typical value of 0.47, see e.g. [10,11]. The same study has discovered that the resistivity behaves as $\rho(T) \propto T^n$ in the high-field non-Fermi liquid (NFL) regime, with $n \simeq 2/3$, while in the low-field NFL regime, at $B \sim B_{c0}$, the exponent $n \simeq 0.45$ [8]. Note that the same behavior of the resistivity was observed in the HF metals URu_2Si_2 [12] and YbAgGe [13] on the verge of the LFL regime, and that the critical behavior takes place up to rather high temperatures comparable with the effective Fermi temperature T_k and up to high magnetic fields. For example, the resistivity measured on $CeCoIn_5$ shows the $T^{2/3}$ behavior over one decade in temperature from 2.3 K to 20 K, and the coefficients $A(B)$ and $C(B)$ exhibit the same behavior at the fields from $B = B_{c0} = 5.1$ T to at least 16 T [8].

In this Letter, we present an explanation of the observed behavior of the electronic system of the heavy-fermion metal $CeCoIn_5$ arisen upon applying magnetic fields B. We analyze the behavior of the effective mass, resistivity, specific heat, charge and heat transport as functions of the applied magnetic field B and show that in the Landau Fermi liquid regime, these quantities demonstrate the critical behavior which is scaled by the critical behavior of the effective mass. In that case, the critical behavior is determined by the fermion condensation quantum phase

transition (FCQPT), whose physics is controlled by quasiparticles with the effective mass which strongly depends on the applied magnetic field B and diverges at $B \to B_{c0}$. In zero magnetic field, the effective mass is controlled by temperature T, and the resistivity is linear in T. In the high-field non-Fermi liquid regime when the system comes from the LFL behavior to the NFL one, the effective mass exhibits very specific behavior, $M^* \sim T^{-2/3}$, and the resistivity demonstrates the $T^{2/3}$ dependence. In the low-field NFL regime, at $B \sim B_{c0}$, this behavior becomes complicated so that the resistivity behaves as T^n, with $n \sim 0.7 - 0.8$. At elevated temperatures and in zero magnetic field, the behavior changes to $M^* \sim T^{-1/2}$, while the resistivity becomes linear in T.

We start with a brief consideration of the LFL regime restored by the application of magnetic field $B > B_{c0}$. If the electronic system approaches FCQPT from the disordered side, the effective mass $M^*(B)$ of the restored LFL depends on magnetic field B as [14,15]

$$M^*(B) \propto \frac{1}{(B - B_{c0})^{2/3}}. \tag{1}$$

Note that Eq. (1) is valid at $T \ll T^*(B)$, where the function $T^*(B) \propto (B - B_{c0})^{4/3}$ determines the line on the $B - T$ phase diagram separating the region of the LFL behavior from the NFL behavior taking place at $T > T^*(B)$ [14]. To estimate the coefficient A, we observe that at the highly correlated regime when $M^*/M \gg 1$, the coefficient $A \propto (M^*)^2$, here M is the bare electron mass [16]. As a result, we have

$$A^2(B) \propto \frac{1}{(B - B_{c0})^{4/3}}, \tag{2}$$

and observe that in the LFL regime, the Kadowaki-Woods ratio, $K = A(B)/\gamma^2(B)$, is conserved because $\gamma(B) \propto M^*(B)$.

Let us now turn to consideration of the system's behavior at elevated temperatures paying special attention to the transition region. To do it, we use the well-know Landau equation relating the quasiparticle energy $\varepsilon(\mathbf{p})$ near the Fermi surface to variations $\delta n(\mathbf{p}, T)$ of the quasiparticle distribution function $n_F(\mathbf{p}, T)$ [17,18]

$$\varepsilon(\mathbf{p}) - \mu = \frac{p_F(p - p_F)}{M^*} + \int F(\mathbf{p}, \mathbf{p}_1)\delta n(\mathbf{p}_1, T)\frac{d\mathbf{p}_1}{(2\pi)^3}. \tag{3}$$

Here, μ is the chemical potential, p_F is the Fermi momentum, $F(\mathbf{p}, \mathbf{p}_1)$ is the Landau amplitude. For the sake of simplicity the summation over the spin variables is omitted. In our case, the variation $\delta n(\mathbf{p}, T)$ is induced by temperature T and defined as $\delta n(\mathbf{p}, T) = n_F(\mathbf{p}, T) - n_F(\mathbf{p}, T = 0)$ with $n_F(\mathbf{p}, T)$ being given by the Fermi-Dirac function

$$n_F(\mathbf{p}, T) = \left\{1 + \exp\left[\frac{(\varepsilon(\mathbf{p}) - \mu)}{T}\right]\right\}^{-1}. \tag{4}$$

Taking into account that $\varepsilon(p \simeq p_F) - \mu = p_F(p - p_F)/M^*$, one directly obtains from Eqs. (4) that $n_F(\mathbf{p}, T \to 0) \to \theta(p_F - p)$, where $\theta(p_F - p)$ is the step function. In our case, Eq. (3) can be used to estimate the behavior of the effective mass $M^*(T)$ as a function of temperature. Actually, differentiating both parts of Eq. (3) with respect to the momentum p, we observe that the difference $p_F/M^*(T) - p_F/M^*(T = 0)$ is given by the integral. In its turn, the integral I

can be estimated upon using the standard procedure of calculating integral when the integrand contains the Fermi-Dirac function, see e.g. [19]. As a result, we obtain that

$$\frac{M}{M^*(T)} \simeq \frac{M}{M^*} + a_1 \left(\frac{TM^*(T)}{T_k M}\right)^2 + a_2 \left(\frac{TM^*(T)}{T_k M}\right)^4 + ... \tag{5}$$

Here a_1 and a_2 are constants proportional to the derivatives of the Landau amplitude with respect to the momentum p. Equation (5) can be regarded as a typical equation of the LFL theory with the only exception for the effective mass M^* which strongly depends on the magnetic field and diverges at $B \to B_{c0}$ as it follows from Eq. (1). Nonetheless, at $T \to 0$, the corrections to $M^*(B)$ start with T^2 terms provided that

$$M/M^*(B) \gg a_1 \left(\frac{TM^*(B)}{T_k M}\right)^2, \tag{6}$$

and the system exhibits the LFL behavior. At some temperature $T_1^*(B) \ll T_k$, the value of the sum on the right hand side of Eq. (5) is determined by the second term. Then Eq. (6) is not valid, and upon omitting the first and third terms, Eq. (5) can be used to determine the effective mass $M^*(T)$ in the transition region,

$$M^*(T) \propto T^{-2/3}. \tag{7}$$

We note, that Eq. (7) has been derived in [15]. Upon comparing Eq. (1) and Eq. (7) and taking into account that the effective mass $M^*(T)$ is a continuous function of T, we can conclude that $T_1^*(B) \propto (B - B_{c0})$.

A few remarks are in order here. Equation (7) is valid if the second term in Eq. (5) is much bigger then the first one, that is

$$\frac{T}{T_k} \gg \left(\frac{M}{M^*}\right)^{3/2}, \tag{8}$$

and this term is bigger then the third one,

$$\frac{T}{T_k} \ll \frac{M}{M^*}. \tag{9}$$

Obviously, both Eq. (8) and (9) can be simultaneously satisfied if $M/M^* \ll 1$. It is seen from Eqs. (1) and (9) that at $B \to B_{c0}$, the range of temperatures over which Eq. (7) is valid shrinks to zero, as well as $T_1^*(B) \to 0$. Thus, it is possible to observe the behavior of the effective mass given by Eq. (7) in a wide range of temperatures provided that the effective mass $M^*(B)$ is diminished by the application of the high magnetic field, see Eq. (1). At $B \to B_{c0}$ and finite temperatures, Eq. (9) cannot be satisfied. Therefore, at elevated temperatures, the third term comes into play making the function $M^*(T)$ be complicated. To estimate the exponent n, we take into account only the third term in Eq. (5) and obtain $M^*(T) \propto T^{-n}$, with $n = 4/5$. As a result, at $B \to B_{c0}$ and $T > T_1^*(B)$, we have an approximation

$$M^*(T) \propto T^{-n}, \tag{10}$$

with the exponent $n \sim 0.7-0.8$. The contribution coming from the other terms can only enlarge the exponent. On the other hand, $n < 1$ because behind FCQPT, when the fermion condensate

is formed, $M^*(T) \propto 1/T$ [20]. Detailed analysis of this item will be published elsewhere. Then, at elevated temperatures, the system comes to a different regime. Smoothing out the step function $\theta(p_F - p)$ at p_F, the temperature creates the variation $\delta n(\mathbf{p}) \sim 1$ over the narrow region $\delta p \sim M^*T/p_F$. In fact, the series on the right hand side of Eq. (5) representing the value of the integral I in Eq. (3) is valid, provided that the interaction radius q_0 in the momentum space of the Landau amplitude F is much larger than δp, $q_0 \gg \delta p$. Otherwise, if $q_0 \sim \delta p$, the series do not represent I and Eqs. (5) and (7) are no longer valid. Such a situation takes place at rising temperatures because the product M^*T grows up as $q_0 \sim \delta p \sim M^*T/p_F \propto T^{1/3}$, as it follows from Eq. (7). As a result, the integral runs over the region q_0 and becomes proportional to M^*T/p_F. Upon omitting the first term on the right hand side of Eq. (3) and substituting the integral by this estimation, we obtain the equation which determines the behavior of the effective mass at $T > T^*(B)$ as [14,21]

$$M^*(T) \propto T^{-1/2}. \tag{11}$$

To capture and summarize the salient features of the LFL behavior observed recently in $CeCoIn_5$ [7,8], we apply the above consideration based on FCQPT. The study of $CeCoIn_5$ in the LFL regime have shown that the coefficients $A(B)$ and $C(B)$, determining the T^2 contributions to the resistivity ρ and thermal resistivity κ_r respectively, possess the same critical field dependence [8]

$$A(B) \propto C(B) \propto \frac{1}{(B - B_{c0})^{4/3}}. \tag{12}$$

The observed critical exponent 4/3 is in excellent agreement with that of given by Eq. (2). Such the parallel behavior of charge and heat transport with the scattering rate growing as T^2 shows that the delocalized fermionic excitations are the Landau quasiparticles carrying charge e. We note that these should be destroyed in the case of conventional quantum phase transitions [2,3]. Nonetheless, let us assume for a moment that these survive. Since the heat and charge transport tend to strongly differ in the presence of the critical fluctuations of superconducting nature, the constancy of the ratio rules out the critical fluctuations [8]. Therefore, we are led to the conclusion that the observed value of the critical magnetic field $B_{c0} = 5.1$ T that coincides approximately with $H_{c2} = 5$ T, the critical field at which the superconductivity vanishes, cannot be considered as giving grounds for the existence of quantum critical behavior of new type. Then, one could expect that some kind of critical fluctuations could cause the observed parallel behavior of charge and heat transport. For example, it is impossible in the case of ferromagnetic fluctuations with a wavevector $q \simeq 0$, but large-q scattering from antiferromagnetic fluctuations of finite momenta could degrade the heat and charge transport in a similar way [11]. In this case, in order to preserve the Kadowaki-Woods ratio these fluctuations are to properly influence the specific heat which characterizes the thermodynamic properties of the system and is not directly related to the transport one. On the other hand, there are no theoretical grounds for the conservation of the Kadowaki-Woods ratio within the frameworks of conventional quantum phase transitions [22]. Therefore, the conservation of the Kadowaki-Woods ratio observed in recent measurements on $CeCoIn_5$ [7] definitely seems to rule out these fluctuations. While both the constancy of Kadowaki-Woods ratio [7] and the constancy of the $A(B)/C(B)$ ratio [8] give strong evidence in favor of the quasiparticle picture.

Now we turn to consideration of the resistivity $\rho(T)$. As we will see below, the striking recent measurements of the resistivity [8,12,13] furnish new evidence in favor of the quasiparticle picture and the existence of FCQPT.

As it follows from Eq. (11) and the mention above relation $A \propto (M^*)^2$, the term $AT^2 \propto M^*T^2$ turns out to be $\propto T$ [14]. As a result, in zero magnetic field and relatively high temperatures $T > T_c$, the resistivity of $CeCoIn_5$ is linear in T. Here T_c is the critical temperature at which the superconductivity vanishes. This observation is in good agreement with experimental facts [5].

At temperatures $T < T_1^*(B)$ and magnetic field $B > B_{c0}$, the system exhibits the LFL behavior with the T^2 dependence of the resistivity $\rho(T)$. Such a behavior is in agreement with experimental facts [6–8].

At the high applied magnetic field and finite temperatures $T > T_1^*(B)$ when the system comes into the NFL regime, the effective mass M^* is determined by Eq. (7). In that case, the range of temperatures over which Eq. (7) is held becomes rather wide, and the system demonstrates the anomalous $T^{2/3}$ resistivity. Actually, upon using the same arguments, we obtain that $AT^2 \propto (M^*)^2T^2 \propto T^{2/3}$ and conclude that the resistivity $\rho(T) \propto T^{2/3}$. Again, this result is in excellent agreement with the reported observations [8,12,13].

If the magnetic field $B \to B_{c0}$ and the temperature is relatively high, $T > T_1^*(B)$, so that the system enters the NFL regime, the effective mass is given by Eq. (10). In that case, the resistivity $\rho(T) \propto (M^*)^2T^2 \propto T^k$, with $k = 2 - 2n = 0.6 - 0.4$. This result is in reasonable agreement with the reported observation of anomalous $T^{0.45}$ dependence of the resistivity in a small region near the critical field $B_{c0} = 5.1$ T [8].

In conclusion, we have shown that the experimentally observed behavior of the electronic system of the heavy-fermion metal $CeCoIn_5$ arisen upon applying magnetic fields can be understood within the frameworks of the FCQPT scenario. We have shown that in the LFL regime the resistivity, specific heat, charge and heat transport as functions of the applied magnetic field B demonstrate the critical behavior which is scaled by the critical behavior of the effective mass. We have observed that this critical behavior is determined by FCQPT, whose physics is controlled by quasiparticles with the effective mass which in the LFL regime strongly depends on the applied magnetic field and diverges at $B \to B_{c0}$. In zero magnetic field, the effective mass is controlled by temperature T, and the resistivity is linear in T. In the high-field NFL regime, the effective mass exhibits very specific behavior, $M^* \sim T^{-2/3}$, while the resistivity demonstrates the $T^{2/3}$ dependence. At elevated temperatures, the behavior changes to $M^* \sim T^{-1/2}$, while the resistivity becomes linear in T.

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